

**Can the Supersymmetric μ parameter be
generated dynamically without a light Singlet?**

RALF HEMPFLING

Deutsches Elektronen-Synchrotron, Notkestraße 85, D-22603 Hamburg, Germany

ABSTRACT

It is generally assumed that the dynamical generation of the Higgs mass parameter of the superpotential, μ , implies the existence of a light singlet at or below the supersymmetry breaking scale, M_{SUSY} . We present a counter-example in which the singlet field can receive an arbitrarily heavy mass (*e.g.*, of the order of the Planck scale, $M_{\text{P}} \approx 10^{19}$ GeV). In this example, a non-zero value of μ is generated through soft supersymmetry breaking parameters and is thus naturally of the order of M_{SUSY} .

The cancellation of quadratic divergences in the unrenormalized Green functions is one of the main motivations of supersymmetry (SUSY). It stabilizes any mass scale under radiative corrections and thus allows the existence of different mass scales such as the electroweak scale, given by the Z boson mass, m_z , and the Planck scale, M_P . The minimal supersymmetric standard model (MSSM) is the most popular model of this kind due to its minimal particle content [1]. In this model, the $SU(2)_L \otimes U(1)_Y$ symmetry breaking is driven by soft SUSY breaking parameters. Thus, the SUSY breaking scale, M_{SUSY} , has to be at or slightly above m_z . For this mechanism to work it is also necessary that the SUSY Higgs mass parameter, $|\mu| \lesssim M_{\text{SUSY}}$. This parameter also determines the chargino and neutralino mass spectrum. From here one can deduce an experimental lower bound from LEP experiments of $|\mu| \gtrsim m_z/4$ independent of $\tan \beta$ [2]. The fact that in the MSSM the μ -parameter, which is a priori arbitrary, has to lie within the narrow range

$$\frac{1}{4}m_z \lesssim |\mu| \lesssim M_{\text{SUSY}}, \quad (1)$$

has been considered a problem of fine-tuning. Possible attempts to try and solve this problem are the inclusion of gravitational couplings [3] or the introduction of additional fields [4].

The introduction of a singlet, N_1 , under the $SU(3)_c \otimes SU(2)_L \otimes U(1)_Y$ standard model (SM) gauge group is the most economical extension of the MSSM in which eq. (1) is natural [4]. Here, the μ parameter is generated dynamically: $\mu = \lambda \langle N_1 \rangle \neq 0$, where $\langle N_1 \rangle$ is the vacuum expectation value (VEV) of N_1 . In this model, one can achieve that the potential vanishes in the direction of N_1 in the SUSY limit by imposing a discrete symmetry. If one includes soft SUSY breaking terms then N_1 acquires a VEV and a mass of order of M_{SUSY} . Thus, one inevitable consequence of this mechanism is the presence of a singlet field under the SM gauge group in the low energy theory. This leads to a severe loss of predictability of the SUSY Higgs sector. In particular, the MSSM prediction

$$m_{h^0} \leq m_z + \text{radiative corrections}, \quad (2)$$

will be evaded [5].

We will demonstrate in the following that it is also possible to make N_1 heavy [say $m_{N_1} = \mathcal{O}(M_P)$] while keeping $\langle N_1 \rangle = \mathcal{O}(M_{\text{SUSY}})$ without fine-tuning. In this limit we recover the predictive Higgs sector of the MSSM [6] with its well defined upper limit of the lightest Higgs boson mass [eq. (2)].

First we need to extend the symmetry group of our Lagrangian in order to forbid the explicit Higgs mass term of the superpotential, $W_H = \mu H \overline{H}$. We choose a continuous symmetry which has to be gauged to avoid a massless Goldstone boson. Our extended gauge group is $\text{SU}(3)_c \otimes \text{SU}(2)_L \otimes \text{U}(1)_Y \otimes \text{U}(1)_{Y'}$. Let us now consider a toy model with three singlets, $N_i \sim (1, 1, 0, Y_i)$ where $Y_i = 2, -2, -1$ for $i = 1, 2, 3$. Here, the first two numbers indicate the multiplicity of N_i under $\text{SU}(3)_c$ and $\text{SU}(2)_L$, and the third and fourth number denote the charges under $\text{U}(1)_Y$ and $\text{U}(1)_{Y'}$. The superpotential of this model is

$$W_N = m N_1 N_2 - \lambda N_1 N_3^2. \quad (3)$$

From here we can derive the SUSY potential, $V_{\text{SUSY}} = V_F + V_{D'}$, where

$$\begin{aligned} V_F &= |m N_2 - \lambda N_3^2|^2 + |m N_1|^2 + 4 |\lambda N_1 N_3|^2, \\ V_{D'} &= \frac{g'^2}{8} (\xi + 2 N_1^* N_1 - 2 N_2^* N_2 - N_3^* N_3)^2. \end{aligned} \quad (4)$$

Here the inclusion of a Fayet-Iliopoulos term [7], ξ , is the easiest way of breaking the $\text{U}(1)_{Y'}$ gauge symmetry but one can envisage other alternatives [8]. The VEVs are denoted by

$$\begin{aligned} n_1 &= \langle N_1 \rangle = 0, \\ n_2 &= \langle N_2 \rangle = \frac{1}{4} \left(-\frac{m}{\lambda} + \sqrt{\frac{m^2}{\lambda^2} + 4\xi} \right), \\ n_3 &= \langle N_3 \rangle = \sqrt{\frac{m n_2}{\lambda}}. \end{aligned} \quad (5)$$

The CP-even and CP-odd components of the scalar field N_1 are mass-degenerate mass-eigenstates with $m_{N_1} = (m^2 + \lambda^2 n_3^2)^{1/2}$. The gauge boson, g' , acquires a mass

$m_{g'} = g'(n_2^2 + n_3^2/4)^{1/2}$ via the Higgs mechanism. The masses of the remaining CP-even (CP-odd) scalars are $m_{N_1}, m_{g'}$ ($m_{N_1}, 0$; the zero mass eigenvalue corresponds to the Goldstone boson which is absorbed to give mass to the gauge boson). The mass eigenvalues of the fermionic components are $\pm m_{N_1}$ and $\pm m_{g'}$ as required if SUSY is unbroken.

Note that in addition to the gauge and the SUSY transformations the Lagrangian is invariant under the global U(1) R -symmetry [9] which does not commute with SUSY. This symmetry transforms $\Phi \rightarrow \exp(in_\Phi \alpha)\Phi$, where $n_\Phi = 2, 0, 0, 0$ for the bosons and $n_\Phi = 1, -1, -1, 1$ for the fermions ($\Phi = N_1, N_2, N_3, g'$). We now break SUSY explicitly in the standard fashion by including soft SUSY breaking terms [10]

$$V_{\text{soft}} = BmN_1N_2 - A\lambda N_1N_3^2 + \text{h.c.}, \quad (6)$$

where $A, B = \mathcal{O}(M_{\text{SUSY}})$ are the soft SUSY breaking parameters. With these terms the R -symmetry is broken down to a discrete Z_2 symmetry ($\alpha = \pm\pi$). If we minimize the full potential, $V = V_{\text{SUSY}} + V_{\text{soft}}$, we find

$$\langle N_1 \rangle \approx (A - B) \frac{mn_2}{m_{N_1}^2} = \mathcal{O}(M_{\text{SUSY}}) \neq 0. \quad (7)$$

We have seen that in our toy-model all the fields acquire a mass of the order of the arbitrary U(1) $_{Y'}$ breaking scale parameterized by $\sqrt{\xi}$. However, the condition $\langle N_1 \rangle = 0$ is protected by R -symmetry to all orders in perturbation theory and is only broken by adding soft SUSY breaking terms [eq. (6)]. We now include in our model the full particle content of the MSSM. The Z_2 symmetry is equivalent to the usual R -parity that prevents baryon and lepton number violating interactions. The full superpotential can then be written as

$$W = W_N + W_H + W_Y, \quad (8)$$

where $W_H = \kappa N_1 H \overline{H}$ and W_Y are the standard Yukawa couplings. The Y' assignments of the quark, leptons and Higgs particles is constrained by the terms

of W in eq. (8) and by requiring the absence of anomalies. These constraints can be satisfied by introducing additional pairs of SUSY multiplets $T \sim (n_c, n_w, Y, Y'_1)$ and $T^c \sim (\bar{n}_c, n_w, -Y, Y'_2)$. These representations have been included in pairs such that below the $U(1)_{Y'}$ breaking scale, $\sqrt{\xi}$, the dynamic mass terms $W \sim m_T T T^c$ can arise ($m_T = \langle N_i \rangle^n / \xi^{(n-1)/2}$; $n = 1, 2, ..$ and $i = 1, 2, 3$). If we assume that all three generations have the same $U(1)_{Y'}$ charges then the absence of anomalies requires the existence of at least one pair of color non-singlets, T and T^c for which $Y'_1 + Y'_2 > 0$ and thus $m_T \propto \langle N_1 \rangle = \mathcal{O}(M_{\text{SUSY}})$. However, the Higgs particles present at M_{SUSY} and thus also the Higgs couplings are equivalent to the MSSM.

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